Subleading corrections to parity violating pion photoproduction

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December 21, 2000

Abstract

We compute the photon asymmetry for near threshold parity violating (PV) pion photoproduction through sub-leading order. We find a potentially significant contribution from new PV low-energy constants not considered previously. We observe that a measurement of this asymmetry may help clarify the theoretical intepretation of backward angle PV electron-proton scattering, provided that the PV πNN Yukawa coupling is determined from nonradiative nuclear PV experiments.

PACS Indices: 11.30.Er, 11.30.Rd, 13.60.Le.

The parity violating (PV) πNN Yukawa coupling constant h_π^1 is a key ingredient to the understanding of the PV nuclear interaction[1, 2, 3, 4, 5]. There have been a number of past attempts to extract this coupling from PV reactions[2, 5, 6, 7, 8, 9], and there are a slate of new experiments, suggested or being planned, with this goal in mind, *e.g.*, $\vec{n}p \to d\gamma$ at LANSCE [10], $\gamma d \to np$ at Jefferson Lab [11], the rotation of polarized neutrons in helium at NIST [9] as well as polarized Compton scattering processes [12, 13].

Recently, Chen and Ji (CJ) suggested using PV π^+ photoproduction near threshold in order to extract h_π^1 [14]. Specifically, the process considered by these authors is

$$\overrightarrow{\gamma}\,(q^\mu;\epsilon^\mu) + p(P_i^\mu) \to \pi^+(k^\mu) + n(P_f^\mu)\ , \qquad (1)$$

where $q^\mu = (\omega, \mathbf{q})$, P_i^μ, $k^\mu = (\omega_\pi, \mathbf{k})$, and P_f^μ are the center-of-mass four-momenta of photon, proton, pion and neutron, respectively, and ϵ^μ is the photon polarization vector. In the threshold region, the pion and photon as well as the nucleon momenta are much smaller than the chiral symmetry breaking scale $\Lambda_\chi = 4\pi F_\pi \sim 1$ GeV and CJ demonstrated that the use of heavy baryon chiral perturbation theory [15, 16] yields a low energy theorem for the threshold PV photon-helicity asymmetry at lowest order of the chiral expansion:

$$A_\gamma\,(\omega_{\rm th}, \theta) = \frac{\sqrt{2}F_\pi(\mu_p - \mu_n)}{g_A m_N} h_\pi^1\ , \qquad (2)$$

The corrections from terms higher order in the chiral expansion were estimated to be around 20% [14]. This process has been studied from the point of view of a more conventional meson exchange approach in [17, 18], and the same h_π^1 dominance of the asymmetry at threshold was found [18].

In this note we present the first order sub-leading result for the aysmmetry. Contrary to the naive expectation that such higher order corrections are small, the subleading correction from the PV πNN vector coupling h_V (defined below) is of order unity, as previously found in the radiative corrections to h_π^1 [19] and to the nucleon anapole moment [20]. We further argue that the proposed parity violating photoproduction process is more suitable for constraining h_V rather than h_π^1. In contrast, other – nonradiative – processes[9, 10, 11] may be most appropriate for the determination of h_π^1, since vector current conservation implies the contributions from h_V are suppressed in such cases. Finally we observe that the contribution from h_V to the nucleon anapole moment introduces significant theoretical uncertainty into the axial vector radiative correction R_A for backward angle PV ep scattering[20]. The SAMPLE collaboration has recently reported an result for this correction which appears to differ from the theoretical estimate [21]. A new constraint on h_V from PV pion photoproduction would reduce the theoretical uncertainty in R_A and may help resolve the apparent difference between theory and experiment.

The motivation behind the use of heavy baryon chiral perturbation theory (HBCPT) is explained in detail in [14], so we follow the notations of this reference. Since we work in the near-threshold region, we use the so-called "small-scale" expansion [22], *i.e.*, we treat $\omega, \omega_\pi, |k|, m_\pi, \delta = m_\Delta - m_N$, etc. as small quantities and characterize amplitudes by the number of powers of these terms, *e.g.*, we count the term $\omega_\pi / q \cdot k$ as being $\mathcal{O}(p^{-1})$. The photon asymmetry arises from the interference of the parity conserving (PC) and PV amplitudes. In Ref. [14] the asymmetry was truncated at leading order, *i.e.*, $\mathcal{O}(p^0)$. In the present work we include the $\mathcal{O}(p)$ correction, which arises dominantly from the PV vector πNN

couplings. As we show later, the corrections from other low energy constants (LECs) like PV $\gamma NN\pi$ vertices are likely to be small, as are other sub-leading contributions to the PV and PC amplitudes. Moreover, chiral loops contribute to the asymmetry only at $\mathcal{O}(p^2)$ and higher. Hence, our truncation of the chiral expansion of the asymmetry is consistent and complete up to terms of $\mathcal{O}(p)$.

The PC amplitudes which describe the charged photoproduction reaction are defined via

$$\begin{aligned} T^{PC} &= N^\dagger \Big[i\mathcal{A}_1\, \sigma\cdot\epsilon + i\mathcal{A}_2 \sigma\cdot\widehat{\mathbf{q}}\ \ \epsilon\widehat{\mathbf{k}} \\ &\qquad + i\mathcal{A}_3\sigma\cdot\widehat{\mathbf{k}}\ \ \epsilon\cdot\widehat{\mathbf{k}} + \mathcal{A}_4 \epsilon\cdot\ \widehat{\mathbf{q}}\times\widehat{\mathbf{k}}\Big] N\ , \end{aligned} \quad (3)$$

where N is the proton Pauli spinor, σ are the Pauli spin matrices, and $\widehat{\mathbf{q}}$ and $\widehat{\mathbf{k}}$ are the unit vectors in the photon and pion directions respectively. At leading order in HBCPT, we have $\mathcal{A}_1 = eg_A/\sqrt{2}F_\pi$, $\mathcal{A}_2 = \mathcal{A}_1\omega|\mathbf{k}|/q\cdot k$, $\mathcal{A}_3 = -\mathcal{A}_1\mathbf{k}^2/q\cdot k$, and $\mathcal{A}_4 = 0$ [23, 24]. As explained in [14] one also requires the non-vanishing subleading order result for $\mathcal{A}_4$.

$$\begin{aligned} \mathcal{A}_4 &= \frac{eg_A\,|\mathbf{k}|}{2\sqrt{2}F_\pi m_N}\left[\mu_p - \left(\frac{\omega}{\omega_\pi}\right)\mu_n\right] \\ &\quad - \frac{2eg_{\pi N\Delta}G_1\,|\mathbf{k}|}{9\sqrt{2}F_\pi m_N}\left(\frac{\omega}{\omega-\delta} + \frac{\omega}{\omega_\pi+\delta}\right)\ , \end{aligned} \quad (4)$$

where the $\Delta(1232)$ contribution has been included explicitly. Here G_1 is the M1 transition moment connecting the nucleon and delta, and $g_{\pi N\Delta}$ is the π-N-Δ coupling. Note that $\mathcal{A}_{1-3}$ is $\mathcal{O}(p^0)$ while $\mathcal{A}_4$ is $\mathcal{O}(p)$.

To $\mathcal{O}(p)$ in the chiral expansion, the PV $\gamma p \to \pi^+ n$ T-matrix can be written as

$$T^{PV} = N^\dagger\left[\mathcal{F}_1\,\widehat{\mathbf{k}}\cdot\epsilon + i\mathcal{F}_2\sigma\cdot\epsilon\times\ \widehat{\mathbf{q}} + i\mathcal{F}_3\sigma\cdot\epsilon\times\ \widehat{\mathbf{k}}\right] N\ . \quad (5)$$

We have then the asymmetry

$$A_\gamma \sim \{\mathcal{A}_1\mathcal{F}_2 + \frac{\sin^2\theta}{2}[\mathcal{A}_3\mathcal{F}_2 - A_4\mathcal{F}_1 - \mathcal{A}_2\mathcal{F}_3] + \cos\theta\mathcal{A}_1\mathcal{F}_3\}\ \ , \quad (6)$$

where $\theta = \cos^{-1}\hat{\mathbf{q}}\cdot\hat{\mathbf{k}}$. The nominally leading piece from the interference term $\mathcal{A}_{1-3}\mathcal{F}_1$ vanishes if the proton target is unpolarized.

The leading, nonvanishing contributions to A_γ, which occur at $\mathcal{O}(p^0)$, are generated by the $\mathcal{O}(p^0)$ terms in $\mathcal{A}_{1-3}$ interfering with the $\mathcal{O}(p^0)$ terms in $\mathcal{F}_2$, and by the $\mathcal{O}(p)$ term in $\mathcal{A}_4$ interfering with the $\mathcal{O}(p^{-1})$ term in $\mathcal{F}_1$. The leading order PV contributions to $\mathcal{F}_{1,2}$ arise from the insertion of the PV Yukawa πNN vertex in FIG. 1 (a), (c), (d).

$$\mathcal{L}^{PV}_{Yukawa} = -ih^1_\pi \pi^+ p^\dagger n + h.c. + \cdots\ , \quad (7)$$

using the sign convention of Ref. [20, 19]. The radiative corrections to h^1_π were discussed extensively in [19], where it was pointed out that what nuclear PV experiments measure is an effective coupling h^{eff}_π [19], which is a linear combination of LECs $h^1_\pi, h_\Delta, h^{(i)}_A$ etc. The commonly used "best value"—$|h^1_\pi| = 5\times10^{-7}$—quoted in [1] corresponds to a large extent to a simple tree-level estimate without loop corrections. The leading-order contribution from the PV Yukawa couplings to the photoproduction amplitudes are found to be[14]:

$$\mathcal{F}_1 = -\frac{eh^1_\pi\,|\mathbf{k}|}{q\cdot k}\ ,\ \mathcal{F}_2 = -\frac{eh^1_\pi}{2m_N}\left[\mu_p - \left(\frac{\omega}{\omega_\pi}\right)\mu_n\right]\ . \quad (8)$$

where $\mathcal{F}_1, \mathcal{F}_2$ are $\mathcal{O}(p^{-1}), \mathcal{O}(p^0)$ respectively.

Subleading contributions to A_γ are generated by $\mathcal{O}(p)$ and $\mathcal{O}(p^2)$ terms in $\mathcal{A}_{1-3}$ and $\mathcal{A}_4$, respectively, inteferencing with the amplitudes in Eq. (8), and by $\mathcal{O}(p)$ contributions in $\mathcal{F}_{2,3}$ interfering with the $\mathcal{O}(p^0)$ terms in $\mathcal{A}_{1-3}$. The subleading PC contributions have been computed in [24]. We refer to the detailed expressions for these corrections in that work, which we employ in our numerical analysis below. Of greater interest are the $\mathcal{O}(p)$ PV amplitudes involving new LEC's. These contributions are generated by the following two Lagrangians:

$$\mathcal{L}^{PV}_V = -\frac{h_V}{\sqrt{2}F_\pi}\bar{p}\gamma^\mu n D_\mu \pi^+ + h.c + \cdots \quad (9)$$

$$\mathcal{L}^{PV}_{\gamma NN\pi} = -ie\frac{C}{\Lambda_\chi F_\pi}\bar{p}\sigma^{\mu\nu}F_{\mu\nu}n\pi^+ + h.c. \quad (10)$$

where $h_V = h^0_V + \frac{4}{3}h^2_V$ in terms of the couplings defined in [19]. We note that the expressions in Eqs. (9, 10) represent only the one-pion term in an expansion of the nonlinear PV Lagrangians of Refs. [4, 19].

These new PV vertices contribute to both the pole diagrams FIG. 1 (c), (d) and the seagull diagram FIG. 1 (b). From the former we have

$$\begin{aligned} \mathcal{F}^{pole}_1 &= -\frac{eh_V\,|\mathbf{k}|}{\sqrt{2}F_\pi m_N}\ , \\ \mathcal{F}^{pole}_2 &= -\frac{eh_V}{2m_N}\frac{\omega_\pi}{\sqrt{2}F_\pi}\left[\mu_p - \left(\frac{\omega}{\omega_\pi}\right)\mu_n\right]\ . \end{aligned} \quad (11)$$

¿From the seagull diagram we find

$$\begin{aligned} \mathcal{F}^{seagull}_1 &= -\frac{eh_V\,|\mathbf{k}|}{2\sqrt{2}F_\pi m_N}\ , \\ \mathcal{F}^{seagull}_2 &= -e\left[\frac{h_V}{2\sqrt{2}m_N} + \frac{2C}{\Lambda_\chi}\right]\frac{\omega}{F_\pi}\ , \\ \mathcal{F}^{seagull}_3 &= \frac{eh_V}{2m_N}\frac{|\mathbf{k}|}{\sqrt{2}F_\pi}\ . \end{aligned} \quad (12)$$

Note that the h_V contribution to $\mathcal{F}_1$ is always two orders higher than that from h^1_π. Thus, since we truncate the asymmetry at $\mathcal{O}(p)$, the terms $\mathcal{F}^{pole}_1, \mathcal{F}^{seagull}_1$ can be safely discarded.

In addition to the $\mathcal{O}(p)$ contributions from h_V and C, $\mathcal{F}_2$ receives an $\mathcal{O}(p)$ contribution involving h^1_π generated by the $1/m_N$ corrections to the nucleon propagator and γNN vertex in the pole amplitudes. We include these corrections in the asymmtry formulae below. Other possible contributions to the PV amplitudes include PV $\gamma N\Delta$ and $\pi N\Delta$ interactions. However, both contribute at $\mathcal{O}(p^2)$, which is higher than the order at which we are truncating. Similarly chiral loop contributions to $\mathcal{A}_{1-4}, \mathcal{F}_1, \mathcal{F}_2, \mathcal{F}_3$ appear at $\mathcal{O}(p^2)$, $\mathcal{O}(p)$, $\mathcal{O}(p^2)$ or higher, respectively. Consequently, chiral loops do not contribute to the asymmetry until at least $\mathcal{O}(p^2)$. All such contributions are then higher order and can be dropped.

Thus the photon asymmetry at order $\mathcal{O}(p)$ reads

$$\begin{aligned}A_\gamma(\omega,\theta) = & \frac{\sqrt{2}h_\pi^1 F_\pi}{g_A m_N \mathcal{G}}\left\{\left[\mu_p - \left(\frac{\omega}{\omega_\pi}\right)\mu_n\right]\left(1-\frac{\sin^2\theta\mathbf{k}^2}{q\cdot k}\right)\right.\\ &\left.+\frac{2}{9}\frac{g_{\pi N\Delta}G_1\sin^2\theta\mathbf{k}^2}{g_A q\cdot k}\left(\frac{\omega}{\omega-\delta}+\frac{\omega}{\omega_\pi+\delta}\right)\right\}\\ &+\frac{h_V\omega_\pi}{g_A m_N\mathcal{G}}\left[\mu_p-\left(\frac{\omega}{\omega_\pi}\right)\mu_n+\frac{\omega}{\omega_\pi}\right]\left(1-\frac{\sin^2\theta\mathbf{k}^2}{2q\cdot k}\right)\\ &+\frac{4\sqrt{2}C\omega}{g_A\Lambda_\chi\mathcal{G}}\left(1-\frac{\sin^2\theta\mathbf{k}^2}{2q\cdot k}\right)\\ &-\frac{|\mathbf{k}|}{g_A m_N\mathcal{G}}[h_V-\frac{2\sqrt{2}F_\pi\omega}{m_N\omega_\pi}\mu_n h_\pi^1]\left(\cos\theta-\frac{\sin^2\theta\omega|\mathbf{k}|}{2q\cdot k}\right)\\ &-\frac{h_\pi^1F_\pi\omega}{\sqrt{2}g_A m_N^2\mathcal{G}}\left[\mu_p-\left(\frac{|\mathbf{k}|^2}{\omega_\pi^2}\right)\mu_n\right]\left(1-\frac{\sin^2\theta\mathbf{k}^2}{2q\cdot k}\right)+\cdots\end{aligned}\tag{13}$$

where $\mathcal{G}=1-\frac{\sin^2\theta|\mathbf{k}|^2}{q\cdot k}\left[1-\frac{(\mathbf{q}-\mathbf{k})^2}{2q\cdot k}\right]$ and where the ellipses indicate the subleading PC contributions of [24].

At threshold—$|\mathbf{k}|=\mathbf{0}$—we find the low energy theorem for the photon asymmetry at $\mathcal{O}(p)$:

$$\begin{aligned}A_\gamma(\omega_{\rm th},\theta) = & \frac{\sqrt{2}F_\pi}{g_A m_N}\left[\mu_p-\mu_n(1+\frac{m_\pi}{2m_N})\right]h_\pi^1\\ &+\frac{m_\pi(\mu_p-\mu_n+1)}{g_A m_N}h_V+\frac{4\sqrt{2}m_\pi}{g_A\Lambda_\chi}C\end{aligned}\tag{14}$$

We observe that the original low energy theorem, Eq. (2), receives a nearly 100% correction from the PV vector coupling h_V, since h_V and h_π^1 may be comparable[4, 20].

There thus exist three separate unknown LEC's in the asymmetry formula, each with its distinct energy dependence and angular distribution, as indicated in FIG. 2. In principle, one might be able to separate h_π^1, h_V, and C by exploiting these different kinematic dependences. In practice, however, this task is probably not feasible given realistic beam time considerations. Thus, we turn to model estimates for guidance as to the magnitudes of the various LEC's. Theoretical values for h_π^1 and h_V have been discussed elsewhere, and one finds no strong *a priori* reason to conclude that they differ significantly in magnitude[20]. For the magnetic LEC C we are not able to perform a first principles theoretical calculation of its size. We may, however, perform an estimate using a vector meson dominance picture involving t-channle ρ- exchange, wherein the parity violation arises from the PV ρNN interaction. When the ρ meson is treated as a heavy degree of freedom and is integrated out, we find the contact effective Lagrangian Eq. (10). The relevant diagram is shown in FIG. 3. For the $\rho\pi\gamma$ vertex we use the Lagrangian:

$$\mathcal{L}^{PC}_{\rho\pi\gamma}=e\frac{g_{\rho\pi\gamma}}{4m_\rho}\epsilon^{\mu\nu\alpha\beta}F_{\mu\nu}G^-_{\alpha\beta}\pi^+ +\cdots\tag{15}$$

where $G_{\alpha\beta}=\partial_\alpha\rho_\beta-\partial_\beta\rho_\alpha$. ¿From the ρ radiative decay width [25] we have $|g_{\rho\pi\gamma}|=0.6$, and for the PV ρNN interaction we follow Ref. [1]

$$\mathcal{L}^{PV}_{\rho NN}=\sqrt{2}(h_\rho^0-\frac{h_\rho^2}{2\sqrt{6}})[\bar{p}\gamma_\mu\gamma_5\rho^+n+H.c.+\cdots]\tag{16}$$

Invoking VMD we have

$$C\sim\frac{g_{\rho\pi\gamma}}{\sqrt{2}}\frac{\Lambda_\chi F_\pi m_\pi}{m_\rho^3}(h_\rho^0-\frac{h_\rho^2}{2\sqrt{6}})\quad .\tag{17}$$

Numerically, the resulting ratio of C and h_π^1 contributions to the asymmetry is

$$\sim 0.017|\frac{h_\rho^0-\frac{h_\rho^2}{2\sqrt{6}}}{h_\pi^1}|=0.03\tag{18}$$

where we have used the "best values" $h_\rho^0=-30g_\pi, h_\rho^2=-25g_\pi, h_\pi=12g_\pi$ [1]. Thus, within the accuracy of our truncation, it appears that one may safely neglect the contribution generated by the interaction in Eq. (10).

One can obtain an experimental constraint on h_π^1 from analysis of a nonradiative PV process. Up to now, however, there exists no such constraint on the magnitude of h_V. A determination of h_V would have both intrinsic interest as well as practical consequences for the interpretation of other experiments. As noted in Ref. [20], the presence of h_V in the nucleon anapole moment – coupled with the present lack of an experimental constraints on this constant – generates considerable theoretical uncertainty in the axial vector radiative corrections to PV ep scattering. Indeed, the numerical coefficient of h_V in R_A is twice as large as the coefficient of the PV ρNN couplings. Given the constraints on the latter from PV $\vec{p}p$ scattering, the PV ρNN couplings give rise to the largest known hadronic effect. It is possible, however, that the h_V contribution is of similar strength. A determination of this LEC from PV pion photoproduction would undoubtedly aid in the intepretation of the SAMPLE result. More generally, the surprising importance of the subleading contribution to A_γ involving h_V suggests careful scrutiny of subleading chiral corrections to other PV radiative processes. The deeper reasons for the importance of subleading chiral corrections in this context as well as others involving hadronic PV [20, 19] is a subject for further study.

Acknowledgment

This work was supported in part under U.S. Department of Energy contracts #DE-AC05-84ER40150 and #DE-FG02-00ER41146, the National Science Foundation, and a National Science Foundation Young Investigator Award.

References

[1] B. Desplanques, J. F. Donoghue and B. R. Holstein, Ann. Phys. 124, 449 (1980).

[2] E. G. Adelberger and W. C. Haxton, Ann. Rev. Nucl. Part. Sci. 35, 501 (1985).

[3] G. B. Feldman et al., Phys. Rev. C 43, 863 (1991).

[4] D. B. Kaplan and M. J. Savage, Nucl. Phys. A 556, 653 (1993).

[5] W. Haeberli and B.R. Holstein in *Symmetries and Fundamental Interactions in Nuclei*, W.C. Haxton and E.M. Henley, Eds., World Scientific, Singapore, 1995, p. 17.

[6] W.T.H. van Oers, Int. J. Mod. Phys. E **8**, 417 (1999).

[7] S.A. Page *et al.*, Phys. Rev. C **35**, 1119 (1987); M. Bini, T. F. Fazzini, G. Poggi, and N. Taccetti, Phys. Rev. C **38**, 1195 (1988).

[8] C.S. Wood et al., Science **275**, 1759 (1997); W.C. Haxton, Science **275**, 1753 (1997); V.V. Flambaum and D.W. Murray, Phys. Rev. C **56**, 1641 (1997); W.S. Wilburn and J.D. Bowman, Phys. Rev. C **57**, 3425 (1998).

[9] V.A. Knyazkov et al., Nucl. Phys. A **417**, 209 (1984); J.F. Cavagnac, B. Vignon and R. Wilson, Phys. Lett. B **67**, 148 (1997); D.M. Markoff, Ph.D. Thesis, University of Washington (1997).

[10] W.M. Snow et al., Nucl. Inst. Meth. A **440**, 729 (2000).

[11] JLab LOI 00-002, W. van Oers and B. Wojtsekhowski, spokesmen.

[12] P.F. Bedaque and M.J. Savage, Phys. Rev. C **62**, 018501 (2000).

[13] J.W. Chen, T.D. Cohen and C.W. Kao, nucl-th/0009031.

[14] J.-W. Chen and X. Ji, hep-ph/0011230.

[15] E. Jenkins and A. V. Manohar, Phys. Lett. B 255, 558 (1991); B 259, 353 (1991).

[16] Ulf-G. Meissner, Int. J. Mod. Phys. E 1, 561 (1992).

[17] R.M. Woloshyn, Can. J. Phys. **57**, 809 (1979).

[18] S.P. Li, E.M. Henley and W-Y. P. Hwang, Ann. Phys. **143**, 372 (1982).

[19] Shi-Lin Zhu, S. Puglia, M. J. Ramsey-Musolf, and B. Holstein, hep-ph/0005281, Phys. Rev. D 63 (in press).

[20] Shi-Lin Zhu, S. Puglia, M. J. Ramsey-Musolf, and B. Holstein, Phys. Rev. D 62, 033008 (2000).

[21] R. Hasty et al, Science 290, 2117 (2000).

[22] T. R. Hemmert, B. R. Holstein and J. Kambor, J. Phys. G 24, 1831 (1998).

[23] V. Bernard, N. Kaiser and U.-G. Meissner, Phys. Lett. B 383, 116 (1996).

[24] H.W. Fearing, T.R. Hemmert, R. Lewis, C. Unkmeir, hep-ph/0005213.

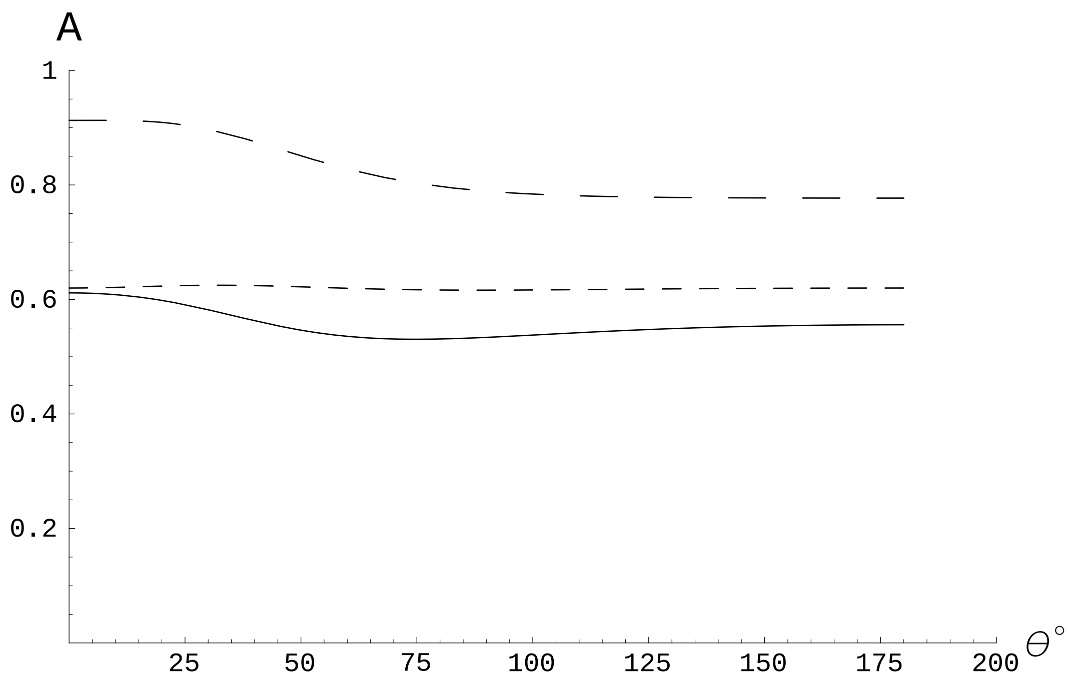
[25] E. Caso et al., Particle Data Group, Euro. Phys. J. C 3, 1 (1998).

Figure Captions

Figure 1. The relevant Feyman diagrams for PV π^+ photoproduction. The circle filled with a cross is the PV vertex.

Figure 2. The θ-dependence of $A_\gamma(\omega = 160 \text{ MeV})$ associated with various PV sources. The solid, long-dashed, and short-dashed curves give kinematic dependences of the coefficients of h_π^1, h_V, and C, respectively.

Figure 3. The t-channel ρ-meson exchange diagram used to estimate the PV LEC C.



A
1
0.8
0.6
0.4
0.2
25
50
75
100
125
150
175
200
θ°